Opening the window to the cogenesis with Affleck-Dine mechanism in gravity mediation

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(Dated: November 30, 2012)

The observed baryon and dark matter densities are equal up to a factor of 5. This observation indicates that the baryon asymmetry and dark matter have the same origin. The Affleck-Dine baryogenesis is one of the most promising mechanisms in this context. Q balls, which are often formed in the early Universe associated with the Affleck-Dine baryogenesis, decay both into supersymmetric particles and into quarks. Recently, it was pointed out that annihilation of squarks into quarks gives a dominant contribution to the Q-ball decay rate and the branching ratio of Q-ball decay into supersymmetric particles changes from the previous estimate. In this paper, the scenario of baryon and dark matter cogenesis from Q ball in gravity mediation is revisited in respect of the improved Q-ball decay rates. It is found that the successful cogenesis takes place when a wino with mass $400 - 600$ GeV is dark matter.

I. INTRODUCTION

The existence of the baryon asymmetry and the dark matter is a long standing challenge in cosmology and particle physics. In supersymmetric (SUSY) extensions of the Standard Model (SM), the lightest SUSY particle (LSP) is a good candidate for dark matter if the R-parity is conserved. Furthermore, the Affleck-Dine mechanism can provide the baryon asymmetry [1, 2]. In the gravity-mediated SUSY breaking model, the Affleck-Dine mechanism often predicts the formation of Q balls in the early universe [3–7]. The Q ball is a spherical condensate of scalar fields. It generally consists of squarks and sleptons, and eventually decays both into quarks and into SUSY particles before the Big Bang Nucleosynthesis (BBN), and the observed baryon asymmetry is released. Through the cascade decays, the SUSY particles produced by the Q-ball decay turn into LSPs, which can account for the dark matter in the Universe. In this case, the baryon asymmetry and dark matter have the same origin and the resultant ratio of baryon to dark matter can be $O(1)$ naturally [4, 8–12].

In the previous works, the branching ratio of the Q-ball decay into SUSY particles is believed to be comparable with that into quarks [4, 9–14]. In this case, the mass of dark matter should be $O(1)$GeV. However, it was pointed out that the many body processes like the squark annihilation may be dominant and then the branching ratio may change drastically [15]. In this letter, we reexamine the branching ratio into SUSY particles in respect of the many body process.

Since the effective mass of the squark inside the Q ball is smaller than that of the free squark, the Q ball can not decay into squarks. We assume that the Q ball is kinematically allowed to decay into binos, winos (LSPs), and SM particles. When the Q-ball decay rate is saturated due to the Pauli exclusion principle [16], the branching ratio is determined only by the number of degrees of

freedom in the final state. Finally, we show that the branching ratio into SUSY particles can be $O(0.01)$. By using this branching ratio, we provide a successful scenario of the baryon and dark matter cogenesis through the Q-ball decay, and show that the wino LSP with mass of $400 - 600$ GeV can naturally explains the observed baryon to dark matter ratio.

This letter is organized as follows. In Sec. II, we briefly review the property of Q balls in gravity mediation. In Sec. III, first we compare the saturated decay and annihilations and then derive the branching ratios. In Sec. IV, we discuss the thermal history in our scenario. Sec. V is devoted to the conclusion.

II. Q BALL PROPERTIES IN GRAVITY MEDIATION

In SUSY extensions of the standard model, there are many flat directions in the scalar potential. The flat directions are lifted by the SUSY breaking effect, and we can take the following potential for the flat direction to see the property of the Q ball in gravity mediation:

$$V = m_\phi^2 |\phi|^2 \left(1 + K \log \frac{|\phi|^2}{M_{\rm P}^2} \right), \qquad (1)$$

where m_ϕ is the mass of the flat direction and $M_{\rm P}$ is the reduced Planck mass($\simeq 2.4 \times 10^{18}$ GeV). In gravity mediation, m_s is the same order of the gravitino mass $m_{3/2}$. The second term in the parenthesis comes from the one-loop radiative corrections, and typically $|K| \sim$0.01-0.1. In many cases, the gluino loops have dominant contributions to the radiative corrections and lead to $K < 0$, and then there exists a Q-ball solution [4, 17]. The energy of the Q ball M_Q, the radius R, the rotation speed of the field ω_0, and the field amplitude at the center of the Q

ball ϕ_0 are given by

$$M_Q \simeq m_\phi(\phi_0)Q, \qquad (2)$$

$$R \simeq \frac{1}{|K|^{1/2}m_\phi(\phi_0)}, \qquad (3)$$

$$\omega_0 \simeq m_\phi(\phi_0), \qquad (4)$$

$$\phi_0 \simeq (2\pi^{3/2})^{-1/2}|K|^{3/4}m_\phi(\phi_0)Q^{1/2}, \qquad (5)$$

where $m_\phi(\phi_0)$ is the mass defined at the energy scale ϕ_0. The rotation speed ω_0 has a further important meaning as $\omega_0 = dM_Q/dQ$; i.e., the Q-ball energy per unit charge.

We are interested in the Q ball with $Q \gtrsim 10^{26}$ as discussed in detail in Sec. IV, and thus the magnitude of the scalar field is $\phi_0 \gtrsim 10^{13}m_\phi(\phi_0)$. At this energy scale, the mass of the flat direction $m_\phi(\phi_0)$ is lower than the mass of squarks at the electro-weak scale due to $K < 0$, and the Q ball can not decay into squarks.

III. Q BALL DECAY RATES INTO BINO-WINO, AND QUARKS

The fermion production rates from the Q ball have upper bounds due to the Pauli exclusion principle [16]. The upper bound of the each massless fermion flux $\boldsymbol{j}$ from the Q-ball surface is calculated as

$$\boldsymbol{n} \cdot \boldsymbol{j} \lesssim 2 \int \frac{d^3k}{(2\pi)^3} \theta\left(\omega_0/2 - |\boldsymbol{k}|\right) \theta(\boldsymbol{k} \cdot \boldsymbol{n}) \hat{\boldsymbol{k}} \cdot \boldsymbol{n}, \quad (6)$$

$$= \frac{2}{8\pi^2} \int_0^{\omega_0/2} k^2 dk = \frac{\omega_0^3}{96\pi^2}, \qquad (7)$$

where $\boldsymbol{n}$ is the outward-pointing normal. We double the flux and take the upper limit of integration as $\omega_0/2$, because one of the decay products has energy less than $\omega_0/2$. We obtain the upper bound for the production rate from the Q ball by multiplying Eq. (7) by the area of the Q-ball surface $4\pi R^2$. The decay rate is saturated when $g\phi_0 > \omega_0$ for the interaction $g\phi\xi\eta$ (ξ, η: massless fermions). The condition $g\phi_0 > \omega_0$ is almost always satisfied due to the large Q value (see Eq. (5)).

In the case of the massive ferimon χ, the upper bound of the flux is lower than Eq. (7). We consider the process of squark $\rightarrow$ quark $+\chi$, and treat the quark as a massless particle. The fermion χ can obtain the energy in the range of $[m_\chi, \omega_0]$, and the quark obtain the energy in the range of $[0, \omega_0 - m_\chi]$. Taking this into account, we just change the integral of Eq. (7) as

$$\frac{1}{8\pi^2} \int_0^{\omega_0 - m_\chi} k^2 dk, \qquad (8)$$

for $\omega_0 > m_\chi > \omega_0/2$, and as

$$\frac{1}{8\pi^2} \left[\int_0^{\omega_0/2} k^2 dk + \int_{m_\chi}^{\omega_0/2} k^2 dk \right], \qquad (9)$$

for $m_\chi < \omega_0/2$. Thus, the χ flux is given by

$$\boldsymbol{n} \cdot \boldsymbol{j}_\chi \simeq \frac{\omega_0^3}{96\pi^2} \times f(m_\chi/\omega_0), \qquad (10)$$

$$f(x) \equiv \begin{cases} 4(1-x)^3 & \text{for } 1/2 < x < 1, \\ 4[(1/2)^3 + (1/2 - x)^3] & \text{for } x < 1/2. \end{cases} \qquad (11)$$

Q balls can also decay into quarks via heavy gluino/higgsino exchange $\phi\phi \rightarrow qq$. This reaction rate is also saturated by the Pauli exclusion principle. The detailed discussion is given in Ref. [7, 18]. The saturated flux is Eq. (7) with ω_0 replaced by $2\omega_0$, which is the total energy available in this process. Thus, we obtain the each quark flux as

$$\boldsymbol{n} \cdot \boldsymbol{j}_{\text{quark}} \simeq \frac{(2\omega_0)^3}{96\pi^2}. \qquad (12)$$

This is larger than Eq. (7) by a factor of 8. Notice that this flux is valid only for $M > \omega_0$, where M is the gluino/higgsino mass, and we assume it in this letter. In Appendix, we show $N(\geq 3)$ body processes are not saturated and negligible.

Now, let us compare the branching ratios of the Q-ball decay into binos, winos, and quarks. The bino or wino production rate is given by Eq. (10), while the quark production rate is given by Eq. (12). Here we should note that since the saturated production rate is determined by the Pauli exclusion principle, the total quark production rate is Eq. (12) times the number of degrees of freedom for quarks produced in the decay. We can count it once we specify the flat direction. Hereafter, we consider the flat direction $\bar{u}_i^a \bar{d}_j^b \bar{d}_k^c \epsilon_{abc}$ ($j \neq k$), where a, b, and c are the color indices and i, j, and k are the family indices. The Q ball can decay into all right handed quarks via gluino exchange and into all left handed quarks via higgsino exchange, because the flat direction contains all colors and, in general, all families. (Even if the flat direction does not contain all families, it can decay into all families through flavor mixings.) The $U(1)_Y$ charge conservation allows one up-type quarks for each two down-type quarks. The Q ball can not directly decay into winos because the $\bar{u}\bar{d}\bar{d}$ flat direction has no tree-level interaction with winos. However, winos are produced via subsequent decays of binos if the LSP is wino. In this letter we consider winos as the LSPs.

We conclude that the total decay rate of the Q ball and the branching ratios of the decay into quarks and bino are calculated as

$$\sum_{\text{all}} \frac{dN}{dt} \simeq \left[8 \times 36 \times \frac{3}{4} + f\left(\frac{m_{\tilde{b}}}{\omega_0}\right) \right] \frac{R^2\omega_0^3}{24\pi}, \quad (13)$$

$$B_{\text{quarks}} \simeq \frac{8 \times 36 \times 3/4}{8 \times 36 \times 3/4 + f(m_{\tilde{b}}/\omega_0)}, \qquad (14)$$

$$B_{\text{bino}} \simeq \frac{f(m_{\tilde{b}}/\omega_0)}{8 \times 36 \times 3/4 + f(m_{\tilde{b}}/\omega_0)}. \qquad (15)$$

In Eq.(13) the factor 36 comes from the degrees of freedom for colors (3), flavors (6) and chiralities (2), and the

factor $3/4$ comes from the $U(1)_Y$ charge conservation. The binos eventually decay into winos (LSPs). Note that the thermal relics of winos do not overclose the Universe.[1]

IV. COGENESIS IN GRAVITY MEDIATION

In this section, we show that our scenario of the baryon and dark matter cogenesis works well. Let us consider the flat direction without non-renormalizable superpotential. The scalar potential for the flat direction is typically written as

$$V(\phi) = (m_\phi^2 - c_H H^2)|\phi|^2 + \frac{m_{3/2}^2}{M_*^{n-2}}(a_m \phi^n + h.c.)$$
$$+ \frac{H^2}{M_*^{n-2}}(a_H \phi^n + h.c.) + \ldots, \qquad (16)$$

where M_* is a cut-off scale and $\ldots$ denotes higher order Planck-suppressed terms. The terms proportional to H^2 are induced via the interaction with the inflaton, and c_H, a_m, and a_H are $O(1)$ constants. Here we assume $c_H > 0$. Owing to the Hubble induced terms and higher order Planck-suppressed terms, the flat direction has a large expectation value during inflation $\phi \simeq M_*$, and then begins to oscillate and rotate around $\phi = 0$ when $H \simeq m_\phi$. Soon after the oscillation, Q balls are formed. Here we assume that the second term in Eq.(16) which kicks ϕ in the phase direction is large enough for ϕ to take a circular orbit. In this case anti-Q balls are not produced, which leads to the simple relation between baryon and dark matter densities.[2] The charge of the Q ball is determined by

$$Q \sim \beta \left(\frac{M_*}{m_\phi}\right)^2 \sim 7 \times 10^{27} \left(\frac{M_*}{M_P}\right)^2 \left(\frac{4\text{TeV}}{m_\phi(\phi_0)}\right)^2, \quad (17)$$

where $\beta = 2 \times 10^{-2}$ [20].

The Q-ball decay temperature is estimated as

$$T_d = \left(\frac{90}{4\pi^2 g_*}\right)^{1/4} \sqrt{\Gamma_Q M_P},$$
$$\simeq 30\text{MeV} \left(\frac{m_\phi(\phi_0)}{4\text{TeV}}\right)^{1/2} \left(\frac{10^{27}}{Q}\right)^{1/2}, \qquad (18)$$

where g_* is the effective relativistic degrees of freedom at the decay time, and $\Gamma_Q = (1/Q) \sum_{\text{all}} dN/dt$ is the decay rate of the Q ball. In the second line of Eq. (18), we set $g_* = 10.75$ and $\sum_{\text{all}} dN/dt \sim 200 \times R^2 \omega_0^3/24\pi Q$. We

find that the Q ball decays before the BBN but after the sphaleron process freezes out [21]. Winos produced from the Q-ball decay do not annihilate when the following condition is satisfied:

$$Y_{\tilde{w}}^{(NT)} \ll \sqrt{\frac{45}{8\pi^2 g_*}} \frac{1}{\langle \sigma v \rangle M_P T_d}, \qquad (19)$$
$$\simeq 3.7 \times 10^{-11} \times \left(\frac{10^{-24}\text{cm}^3/\text{s}}{\langle \sigma v \rangle}\right)\left(\frac{30\,\text{MeV}}{T_d}\right)(20)$$

As mentioned above, we consider the flat direction without non-renormalizable superpotential. In this case, the Q balls dominate the Universe soon after inflation, and the baryon-to-entropy ratio is given by

$$Y_b \simeq 10^{-10} \left(\frac{T_d}{4\text{MeV}}\right)\left(\frac{100\text{TeV}}{m_\phi}\right)\left(\frac{10^2}{\Delta}\right), \qquad (21)$$

where we include the dilution factor Δ. There is some mechanism to produce entropy after the reheating of the inflation, such as thermal inflation [22] and domain wall decay [23]. We do not specify the dilution mechanism and assume that the baryon asymmetry produced from the Q-ball decay is consistent with the observation.

In the case of the wino LSP, the thermal relic abundance can be ignored for $m_{\tilde{w}} \ll 1\text{TeV}$ [24]. The baryon-to-dark matter ratio is determined only by the Q-ball decay:

$$5 \simeq \frac{\Omega_{\text{DM}}}{\Omega_b} = \frac{3m_{\tilde{w}}}{m_N} \frac{B_{\text{bino}}}{B_{\text{quarks}}}. \qquad (22)$$

From Eqs. (11), (15), and (22), the bino and wino masses are related with each other by the following equation:

$$\frac{360\text{GeV}}{m_{\tilde{w}}} = f\left(\frac{m_{\tilde{b}}}{\omega_0}\right), \qquad (23)$$

where $f(x)$ is defined as Eq. (11). The results are shown in Fig. 1 for the case of $m_{\tilde{w}} = m_{\tilde{b}}$. There are two solutions when $\omega_0 \gtrsim 1.5\text{TeV}$. In the limit of $\omega_0 \to \infty$, two solutions are approximated to 360GeV and ω_0. For $3\text{TeV} < m_\phi(\phi_0) < 10\text{TeV}$, the wino mass is predicted to be $400 - 600$ GeV. The resultant wino abundance is given by

$$Y_{\tilde{w}}^{(NT)} \simeq 1.1 \times 10^{-12} \frac{360\text{GeV}}{m_{\tilde{w}}}. \qquad (24)$$

From this and Eq. (20), we can check that the winos do not annihilate. Indirect detection experiments constrain the wino mass as $m_{\tilde{w}} \gtrsim 300\text{GeV}$ [25–27]. The above predicted wino mass satisfies this constraint.

V. CONCLUSIONS

We have reinvestigated the baryon and dark matter cogenesis through Q-ball decay into quarks and SUSY particles by taking into account the squark annihilation

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[1] On the other hand, if binos with mass larger than a few $\times 100$ GeV are the LSPs, their relic density may overcloses the Universe .

[2] When both Q-balls and anti-Q balls are formed, cogenesis requires much smaller LSP mass. Such small LSP mass is realized in gauge mediated SUSY breaking where a gravitino with mass $\lesssim 1$ GeV is the LSP [19].
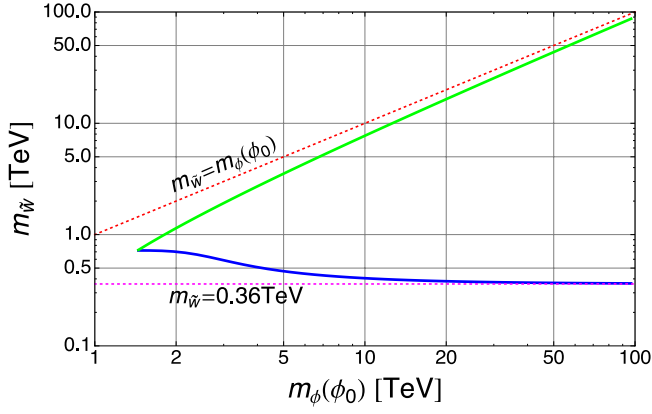
FIG. 1. Solutions of Eq. (23) for the case of $m_{\tilde{w}} = m_{\tilde{b}}$ as a function of $m_\phi(\phi_0)$. There is no solution for $m_\phi(\phi_0) \lesssim 1.5$TeV and are two independent solutions for $m_\phi(\phi_0) \gtrsim 1.5$TeV (green and blue lines). The red and magenta dotted lines show the two asymptotic solutions $m_{\tilde{w}} = m_\phi(\phi_0)$ and $m_{\tilde{w}} = 360$GeV, respectively.

process inside the Q ball. The branching ratio of the Q-ball decay into quarks is enhanced by the number of degrees of freedom for quarks produced in the decay. We have assumed that the Q ball can decay into binos, winos, and SM particles kinematically, and considered the wino as LSP. In this case, we show that the branching into binos can be $O(0.01)$ for the $\bar{u}\bar{d}\bar{d}$ flat direction and predict that the dark matter is the wino with mass of $400 - 600$ GeV.

Appendix: Q-ball decay rates through the $N \geq 3$ body scattering processes

Not only the decay process but also the N body scattering processes can occur in the Q ball. The rate of the charge emission from the Q ball through the N body scattering process can be roughly estimated as

$$\left(\frac{\mathrm{d}N}{\mathrm{d}t}\right)_N \sim Q \times n_\phi^{N-1} \times \Gamma_N, \qquad (A.1)$$

$$\Gamma_N = \int \mathrm{dLips}|\mathcal{M}|^2 \prod_{\text{initial}} \frac{1}{2E_i}, \qquad (A.2)$$

$$\mathrm{dLips} \equiv (2\pi)^4 \delta\left(\sum_{\text{all}} p_j\right) \prod_{\text{final}} \frac{\mathrm{d}^3 k_i}{(2\pi)^3 2E_i}, \quad (A.3)$$

where $n_\phi \sim \omega_0 \phi_0^2$ is the squark number density in the Q ball. Let us show that the rates of the N body scattering processes are not saturated for $N \geq 3$.

The mass of the field interacting with the Q ball is $O(\phi_0)$, but the typical interaction energy is $O(\omega_0)$. Thus, we can estimate the rates of the N body scattering processes in the leading order of $\omega_0/\phi_0 \sim Q^{-1/2}$. The number of particles in the final state should be minimized in

the leading order as

$$N_{\text{ext}} = \begin{cases} N, & \text{N: even} \\ N+1, & \text{N: odd.} \end{cases} \qquad (A.4)$$

Then, the number of fermion propagators can be counted as

$$N_{\text{prop}} = \begin{cases} 3N/2 - 2, & \text{N: even} \\ 3N/2 - 3/2, & \text{N: odd.} \end{cases} \qquad (A.5)$$

However, as shown in Fig. 2, there should be a factor of M from the chirality flip, where M is the Majorana gluino mass or the higgsino mass, and we assume $\omega_0 < M \ll \phi_0$. The number of mass insertions is

$$N_{\text{mass}} = \begin{cases} N/2, & \text{N: even} \\ (N-1)/2, & \text{N: odd.} \end{cases} \qquad (A.6)$$

The gauge boson is massless if it has no tree level interaction with the Q ball. Hereafter, we conservatively take the gauge boson as a massless field. Thus, from Eqs. (A.5) and (A.6), we can estimate $|\mathcal{M}|^2$ as

$$|\mathcal{M}|^2 \sim \begin{cases} \phi_0^{8-6N} M^N \omega_0^N, & \text{N: even} \\ \phi_0^{6-6N} M^{N-1} \omega_0^{N+1}, & \text{N: odd} \end{cases} \qquad (A.7)$$

Here we determine the ω_0 dependence from dimensional analysis. On the other hand, the kinematics is determined only by ω_0. We conclude that the charge emission rates from Q ball through the N body scattering process can be estimated as

$$\left(\frac{dN}{dt}\right)_N \sim Q(\omega_0 \phi_0^2)^{N-1} \Gamma_N,$$

$$\sim \begin{cases} \omega_0 Q^{4-2N}(M/\omega_0)^N, & \text{N: even} \\ \omega_0 Q^{3-2N}(M/\omega_0)^{N-1}, & \text{N: odd} \end{cases} \qquad (A.8)$$

where we have used $Q \sim \phi_0^2/\omega_0^2$. We should compare this with the saturated emission rate from the Q ball $(\mathrm{d}N/\mathrm{d}t)_{\text{sat}} \sim \omega_0$ (see Eqs. (3) and (13)) and find that the rate is not saturated for $N \geq 3$.
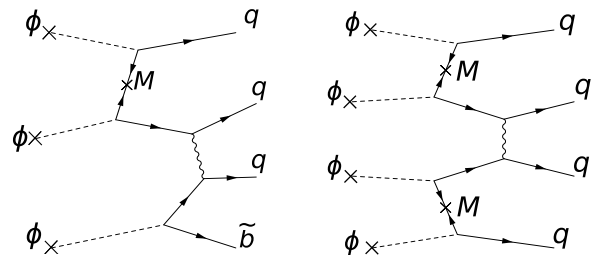

FIG. 2. Examples of the diagrams for the N body scattering processes.

[1] I. Affleck and M. Dine, Nucl. Phys. B **249**, 361 (1985).

[2] M. Dine, L. Randall and S. D. Thomas, Nucl. Phys. B **458**, 291 (1996).

[3] A. Kusenko and M. E. Shaposhnikov, Phys. Lett. B **418**, 46 (1998).

[4] K. Enqvist and J. McDonald, Phys. Lett. B **425**, 309 (1998); Nucl. Phys. B **538**, 321 (1999).

[5] S. Kasuya and M. Kawasaki, Phys. Rev. D **61**, 041301(R) (2000).

[6] S. Kasuya and M. Kawasaki, Phys. Rev. D **62**, 023512 (2000).

[7] S. Kasuya and M. Kawasaki, Phys. Rev. D **64**, 123515 (2001).

[8] E. Komatsu *et al.* [WMAP Collaboration], Astrophys. J. Suppl. **192**, 18 (2011).

[9] D. B. Kaplan, Phys. Rev. Lett. **68**, 741 (1992).

[10] M. Fujii and K. Hamaguchi, Phys. Lett. B **525**, 143 (2002).

[11] M. Fujii and K. Hamaguchi, Phys. Rev. D **66**, 083501 (2002).

[12] M. Fujii and T. Yanagida, Phys. Lett. B **542**, 80 (2002).

[13] I. M. Shoemaker and A. Kusenko, Phys. Rev. D **80**, 075021 (2009).

[14] F. Doddato and J. McDonald, JCAP **1106**, 008 (2011).

[15] S. Kasuya and M. Kawasaki, Phys. Rev. D **84**, 123528 (2011).

[16] A. G. Cohen, S. R. Coleman, H. Georgi and A. Manohar, Nucl. Phys. B **272**, 301 (1986).

[17] S. Coleman, Nucl. Phys. **B262** (1985) 263.

[18] M. Kawasaki and M. Yamada, arXiv:1209.5781 [hep-ph].

[19] S. Kasuya, M. Kawasaki and M. Yamada, arXiv:1211.4743 [hep-ph].

[20] T. Hiramatsu, M. Kawasaki and F. Takahashi, JCAP **1006**, 008 (2010).

[21] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B **155**, 36 (1985).

[22] D. H. Lyth and E. D. Stewart, Phys. Rev. D **53**, 1784 (1996); Phys. Rev. Lett. **75**, 201 (1995).

[23] M. Kawasaki and F. Takahashi, Phys. Lett. B **618**, 1 (2005).

[24] J. Hisano, S. Matsumoto, M. Nagai, O. Saito and M. Senami, Phys. Lett. B **646**, 34 (2007).

[25] A. Charbonnier, C. Combet, M. Daniel, S. Funk, J. A. Hinton, D. Maurin, C. Power and J. I. Read *et al.*, Mon. Not. Roy. Astron. Soc. **418**, 1526 (2011).

[26] O. Adriani *et al.* [PAMELA Collaboration], Phys. Rev. Lett. **105**, 121101 (2010).

[27] M. Ibe, S. Matsumoto and T. T. Yanagida, Phys. Rev. D **85**, 095011 (2012).